

Dressing the inflaton with the Standard Model gauge group

A. Mazumdar^{1, 2}

¹ *Physics Department, Lancaster University, LA1 4YB, United Kingdom*

² *Niels Bohr Institute, Blegdamsvej-17, Copenhagen-2100, Denmark*

ABSTRACT: In this talk we will discuss how inflation can be embedded within a minimal extension of the Standard Model where the inflaton carries the Standard Model charges. There is no need of an ad-hoc scalar field to be introduced in order to explain the temperature anisotropy of the cosmic microwave background radiation, all the ingredients are present within a minimal supersymmetric Standard Model. The model is robust enough to provide a successful exit from inflation with all observed matter in the universe. This is a triumph for an inflationary paradigm which has always begged a simple question: can we identify the inflaton in a laboratory. We will briefly discuss how LHC can shed some insight into the inflaton.

Inflation has been extremely successful in explaining the temperature anisotropy of the observed cosmic microwave background radiation by generating almost scale invariant density perturbations [1]. (for a theory of density perturbations see [2]). It has been known for almost 26 years that inflation can be driven by a dynamical scalar field known as the *inflaton*, an order parameter, which could either be fundamental or composite. Particularly, if the inflaton rolls very slowly on a sufficiently flat potential, such that the potential energy density dominates over the kinetic term, then all the successes of inflation can be met, i.e. dynamical explanation of the homogeneity and the isotropy of the universe on very large scales (for a review see [3]), Gaussian density perturbations, etc.

However, inspite of the impressive list of achievements, inflation was never embedded in a fundamental theory which could also be testable in a laboratory. Inspite of many attempts there has been no single good candidate for an inflaton which comes naturally out of a well motivated theory of particle physics (for a review on models of inflation, see [4]). One always relies on scalar fields which are *absolute gauge singlets* possibly residing in some hidden sector or secluded sector with a small coupling to the SM gauge group. By definition an *absolute gauge singlet* does not carry any charge what so-ever be the case. Therefore the masses, couplings and interactions are not generally tied to any fundamental theory or any symmetry. Such gauge singlets are used ubiquitously by model builders to obtain a desired potential and interactions at a free will in order to explain the current CMB data. Not only that many notable papers use ad-hoc couplings to explain phenomena such as preheating and thermalization without bothering the relevant degrees of freedom required to create a Universe like ours.

Very recently some of these questions have been addressed in a low energy field theory setup which explains (for a review see [5]):

- the origin of inflation
- the fundamental interactions of an inflaton
- how the inflaton creates Standard Model baryons and cold dark matter ?
- and how can we test the inflaton in a laboratory ?

For the first time we are aiming to build a holistic model of inflation which is truly embedded in a Standard Model (SM) gauge theory. The inflaton carries

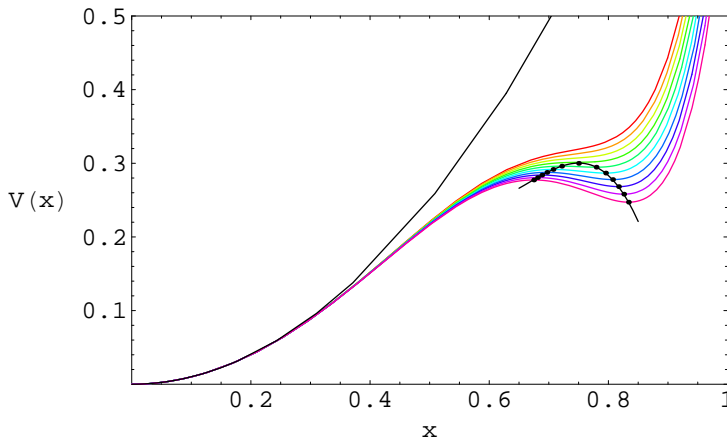


Figure 1: The colored curves depict the full potential, where $V(x) \equiv V(\phi)/(0.5 m_\phi^2 M_P^2 (m_\phi/M_P)^{1/2})$, and $x \equiv (\lambda_n M_P/m_\phi)^{1/4}(\phi/M_P)$. The black curve is the potential arising from the soft SUSY breaking mass term. The black dots on the colored potentials illustrate the gradual transition from minimum to the saddle point and to the maximum.

the SM charges and inflation occurs within an observable sector if the low energy supersymmetry (SUSY) is found just above the electroweak scale at the LHC [6, 7, 8, 9, 10, 11, 12, 13]. We will explain below why SUSY is necessary to build a successful inflation model. SUSY provides the scalar fields (partners of the SM fermions and gauge bosons) and the stability of the flatness of the inflaton potential.

Furthermore, since the inflaton carries the SM charges, it decays *only* into the SM particles and SUSY particles, i.e. quarks, squarks, leptons, sleptons, etc. Within a minimal supersymmetric Standard Model (MSSM) we know all the relativistic species and therefore we can trace back thermal history of the universe accurately above the electroweak scale. We highlight that all the relevant physical processes are happening within an observable sector alone. If the lightest supersymmetric particle (LSP) is stable due to R-parity, we naturally obtain cold dark matter as a consequence of inflation [10].

Let us now concentrate on the model of inflation which is based on MSSM together with gravity [6, 7, 10]. Therefore consistency dictates that all non-renormalizable terms allowed by gauge symmetry and supersymmetry should be included below the cut-off scale, which we take to be the Planck scale. The superpotential term which

lifts the F -flatness is given by (see for a review [14]):

$$W_{non} = \sum_{n>3} \frac{\lambda_n}{n} \frac{\Phi^n}{M^{n-3}}, \quad (1)$$

where Φ is a *gauge invariant* superfield which contains the flat direction. Within MSSM all the flat directions are lifted by non-renormalizable operators with $4 \leq n \leq 9$ [15], where n depends on the flat direction. We expect that quantum gravity effects yield $M = M_P = 2.4 \times 10^{18}$ GeV and $\lambda_n \sim \mathcal{O}(1)$ [16]. Note however that our results will be valid for any values of λ_n , because rescaling λ_n simply shifts the VEV of the flat direction. Let us focus on the lowest order superpotential term in Eq. (1) which lifts the flat direction. Soft SUSY breaking induces a mass term for ϕ and an A -term so that the scalar potential along the flat direction reads

$$V = \frac{1}{2} m_\phi^2 \phi^2 + A \cos(n\theta + \theta_A) \frac{\lambda_n \phi^n}{n M_P^{n-3}} + \lambda_n^2 \frac{\phi^{2(n-1)}}{M_P^{2(n-3)}}, \quad (2)$$

Here ϕ and θ denote respectively the radial and the angular coordinates of the complex scalar field $\Phi = \phi \exp[i\theta]$, while θ_A is the phase of the A -term (thus A is a positive quantity with dimension of mass). Note that the first and third terms in Eq. (2) are positive definite, while the A -term leads to a negative contribution along the directions whenever $\cos(n\theta + \theta_A) < 0$, see 1. The importance of the A -term was first highlighted in a successful MSSM curvaton model [17]¹.

The maximum impact from the A -term is obtained when $\cos(n\theta + \theta_A) = -1$ (which occurs for n values of θ). For a choice of parameter, $A^2 = 8(n-1)m_\phi^2$, the potential is flat near ϕ_0 , where first and second derivative of the potential vanishes, i.e. $V'(\phi_0) = 0$, $V''(\phi_0) = 0$. As a result, if initially $\phi \sim \phi_0$, a slow roll phase of inflation is driven by the third derivative of the potential. We will show that the potential near ϕ_0 becomes cosmologically flat. As a matter of fact in the gravity mediated SUSY breaking case, the A -term and the soft SUSY breaking mass terms are expected to be the same order of magnitude as the gravitino mass, i.e. $m_\phi \sim A \sim m_{3/2} \sim \mathcal{O}(1)$ TeV see [19].

If the above condition is not satisfied then for $A^2 \geq 8(n-1)m_\phi^2$ the potential develops a secondary false minimum with a charge and color breaking minimum. Although inflation could occur there [20] but phenomenologically such a situation is not desired at all. In the other extreme case it is possible to drive assisted [21]

¹Similar ideas for the curvaton to carry the SM charges have been entertained in [18].

inflation with many flat directions, however, such a possibility does not arise within MSSM [22].

The potential near the saddle point, $A^2 = 8(n-1)m_\phi^2$, is very flat along the *real direction* but not along the *imaginary direction*. Along the imaginary direction the curvature is determined by m_ϕ . Around ϕ_0 the field lies in a plateau with a potential energy

$$V(\phi_0) = \frac{(n-2)^2}{2n(n-1)} m_\phi^2 \phi_0^2 \quad (3)$$

with

$$\phi_0 = \left(\frac{m_\phi M_{\text{P}}^{n-3}}{\lambda_n \sqrt{2n-2}} \right)^{1/(n-2)}. \quad (4)$$

This results in Hubble expansion rate during inflation which is given by

$$H_{\text{inf}} = \frac{(n-2)}{\sqrt{6n(n-1)}} \frac{m_\phi \phi_0}{M_{\text{P}}}. \quad (5)$$

When ϕ is very close to ϕ_0 , the first derivative is extremely small. The field is effectively in a de Sitter background, and we are in self-reproduction (or *eternal inflation*) regime where the two point correlation function for the flat direction fluctuation grows with time. But eventually classical friction wins and slow roll begins at $\phi \approx \phi_{\text{self}}$ [6, 7]

$$(\phi_0 - \phi_{\text{self}}) \simeq \left(\frac{m_\phi \phi_0^2}{M_{\text{P}}^3} \right)^{1/2} \phi_0. \quad (6)$$

The regime of *eternal inflation* plays an important role in addressing the initial condition problem [11].

The observationally relevant perturbations are generated when $\phi \approx \phi_{\text{COBE}}$. The number of e-foldings between ϕ_{COBE} and ϕ_{end} , denoted by $\mathcal{N}_{\text{COBE}}$:

$$\mathcal{N}_{\text{COBE}} \simeq \frac{\phi_0^3}{2n(n-1)M_{\text{P}}^2(\phi_0 - \phi_{\text{COBE}})}. \quad (7)$$

The amplitude of perturbations thus produced is given by [7]

$$\delta_H \equiv \frac{1}{5\pi} \frac{H_{\text{inf}}^2}{\dot{\phi}} \simeq \frac{1}{5\pi} \sqrt{\frac{2}{3} n(n-1)(n-2)} \left(\frac{m_\phi M_{\text{P}}}{\phi_0^2} \right) \mathcal{N}_{\text{COBE}}^2, \quad (8)$$

and the spectral tilt of the power spectrum and its running are found to be [6, 7]

$$n_s = 1 + 2\eta - 6\epsilon \simeq 1 - \frac{4}{\mathcal{N}_{\text{COBE}}}, \quad (9)$$

$$\frac{d n_s}{d \ln k} = -\frac{4}{\mathcal{N}_{\text{COBE}}^2}. \quad (10)$$

As discussed in [6, 7], among nearly 300 flat directions there are two that can lead to a successful inflation along the lines discussed above.

One is udd which, up to an overall phase factor, is parameterized by

$$u_i^\alpha = \frac{1}{\sqrt{3}}\phi, \quad d_j^\beta = \frac{1}{\sqrt{3}}\phi, \quad d_k^\gamma = \frac{1}{\sqrt{3}}\phi. \quad (11)$$

Here $1 \leq \alpha, \beta, \gamma \leq 3$ are color indices, and $1 \leq i, j, k \leq 3$ denote the quark families. The flatness constraints require that $\alpha \neq \beta \neq \gamma$ and $j \neq k$.

The other direction is LLe ², parameterized by (again up to an overall phase factor)

$$L_i^a = \frac{1}{\sqrt{3}} \begin{pmatrix} 0 \\ \phi \end{pmatrix}, \quad L_j^b = \frac{1}{\sqrt{3}} \begin{pmatrix} \phi \\ 0 \end{pmatrix}, \quad e_k = \frac{1}{\sqrt{3}}\phi, \quad (12)$$

where $1 \leq a, b \leq 2$ are the weak isospin indices and $1 \leq i, j, k \leq 3$ denote the lepton families. The flatness constraints require that $a \neq b$ and $i \neq j \neq k$. Both these flat directions are lifted by $n = 6$ non-renormalizable operators [15],

$$W_6 \supset \frac{1}{M_{\text{P}}^3}(LLe)(LLe), \quad W_6 \supset \frac{1}{M_{\text{P}}^3}(udd)(udd). \quad (13)$$

The reason for choosing either of these two flat directions³ is twofold: (i) a non-trivial A -term arises, at the lowest order, only at $n = 6$; and (ii) we wish to obtain the correct COBE normalization of the CMB spectrum.

Those MSSM flat directions which are lifted by operators with dimension $n = 7, 9$ are such that the superpotential term contains at least two monomials, i.e. is of the type

$$W \sim \frac{1}{M_{\text{P}}^{n-3}} \Psi \Phi^{n-1}. \quad (14)$$

If ϕ represents the flat direction, then its VEV induces a large effective mass term for ψ , through Yukawa couplings, so that $\langle \psi \rangle = 0$. Hence Eq. (14) does not contribute to the A -term.

More importantly, as we will see, all other flat directions except those lifted by $n = 6$ fail to yield the right amplitude for the density perturbations. Indeed, as can be seen in Eq. (4), the value of ϕ_0 , and hence also the energy density, depend on n .

²When the flat direction develops a VEV during inflation, it spontaneously breaks $SU(2) \times U(1)_y$, which gives masses to the corresponding gauge bosons. It is possible to obtain a seed perturbations for the primordial magnetic field in this case, see [23].

³Since LLe and udd are independently D - and F -flat, inflation could take place along any of them but also, at least in principle, simultaneously. The dynamics of multiple flat directions are however quite involved [24].

According to the arguments presented above, successful MSSM flat direction inflation has the following model parameters:

$$m_\phi \sim 1 - 10 \text{ TeV}, \quad n = 6, \quad A = \sqrt{40}m_\phi, \quad \lambda \sim \mathcal{O}(1). \quad (15)$$

Here we assume that λ (we drop the subscript "6") is of order one, which is the most natural assumption when $M = M_{\text{P}}$.

The Hubble expansion rate during inflation and the VEV of the saddle point are ⁴

$$H_{\text{inf}} \sim 1 - 10 \text{ GeV}, \quad \phi_0 \sim (1 - 3) \times 10^{14} \text{ GeV}. \quad (16)$$

Note that both the scales are sub-Planckian. The total energy density stored in the inflaton potential is $V_0 \sim 10^{36} - 10^{38} \text{ GeV}^4$. The fact that ϕ_0 is sub-Planckian guarantees that the inflationary potential is free from the uncertainties about physics at super-Planckian VEVs. The total number of e-foldings during the slow roll evolution is large enough to dilute any dangerous relic away [7],

$$\mathcal{N}_{\text{tot}} \sim 10^3, \quad (17)$$

At such low scales as in MSSM inflation the number of e-foldings, $\mathcal{N}_{\text{COBE}}$, required for the observationally relevant perturbations, is much less than 60 [25]. If the inflaton decays immediately after the end of inflation, we obtain $\mathcal{N}_{\text{COBE}} \sim 50$. Despite the low scale, the flat direction can generate adequate density perturbations as required to explain the COBE normalization. This is due to the extreme flatness of the potential (recall that $V' = 0$), which causes the velocity of the rolling flat direction to be extremely small. From Eq. (8) we find an amplitude of

$$\delta_H \simeq 1.91 \times 10^{-5}. \quad (18)$$

There is a constraint on the mass of the flat direction from the amplitude of the CMB anisotropy:

$$m_\phi \simeq (100 \text{ GeV}) \times \lambda^{-1} \left(\frac{\mathcal{N}_{\text{COBE}}}{50} \right)^{-4}. \quad (19)$$

We get a lower limit on the mass parameter when $\lambda \leq 1$. For smaller values of $\lambda \ll 1$, the mass of the flat direction must be larger. Note that the above bound on the inflaton mass arises at high scales, i.e. $\phi = \phi_0$. However, through renormalization group flow, it is connected to the low scale mass, as will be discussed in Sect. 4.

⁴We note that H_{inf} and ϕ_0 depend very mildly on λ as they are both $\propto \lambda^{-1/4}$.

The spectral tilt of the power spectrum is not negligible because, although the first slow roll parameter is $\epsilon \sim 1/\mathcal{N}_{\text{COBE}}^4 \ll 1$, the other slow roll parameter is given by $\eta = -2/\mathcal{N}_{\text{COBE}}$ and thus, see Eq. (9)⁵

$$n_s \sim 0.92, \quad (20)$$

$$\frac{dn_s}{d \ln k} \sim -0.002, \quad (21)$$

where we have taken $\mathcal{N}_{\text{COBE}} \sim 50$ (which is the maximum value allowed for the scale of inflation in our model). In the absence of tensor modes, this agrees with the current WMAP 3-years' data within 2σ [1]. Note that MSSM inflation does not produce any large stochastic gravitational wave background during inflation. Gravity waves depend on the Hubble expansion rate, and in our case the energy density stored in MSSM inflation is very small. Inflation can still happen for small deviations from the saddle point condition, $A^2 = 8(n-1)m_\phi^2$. Notable point is that the spectral tilt can match the current observations, i.e. within $0.92 \leq n_s \leq 1.0$. The plot above summarizes the results.

The blue band above shows the experimentally allowed region. We find that smaller values of m_ϕ are preferred for smaller values of n_s . We also find that the allowed range of m_ϕ is $75 - 440$ GeV for the experimental ranges of n_s and δ_H . We assume $\lambda \sim 1$ for these two figures. If λ is less than $\mathcal{O}(1)$, e.g., $\lambda \sim 0.1$ or so (which can occur in $SO(10)$ model), it will lead to an increase in m_ϕ .

Since the MSSM inflaton candidates are represented by *gauge invariant* combinations which are not singlets. The inflaton parameters receive corrections from gauge interactions which, unlike in models with a gauge singlet inflaton, can be computed in a straightforward way. Quantum corrections result in a logarithmic running of the soft supersymmetry breaking parameters m_ϕ and A . The conclusion is robust, although the soft terms and the value of the saddle point are all affected by radiative corrections, they do not remove the saddle point nor shift it to unreasonable values. The existence of a saddle point is thus insensitive to radiative corrections.

One final comment is in order before closing this Section. Unlike m_ϕ , there is no prospect of measuring the A term, because it is related to the non-renormalizable interactions which are suppressed by M_{P} . However, a knowledge of supersymmetry

⁵Obtaining $n_s > 0.92$ (or $n_s < 0.92$, which is however outside the 2σ allowed region) requires deviation from the saddle point condition, $A^2 = 8(n-1)m_\phi^2$, see the discussion below. For a more detailed discussion on the spectral tilt, see also Refs. [26],[12].

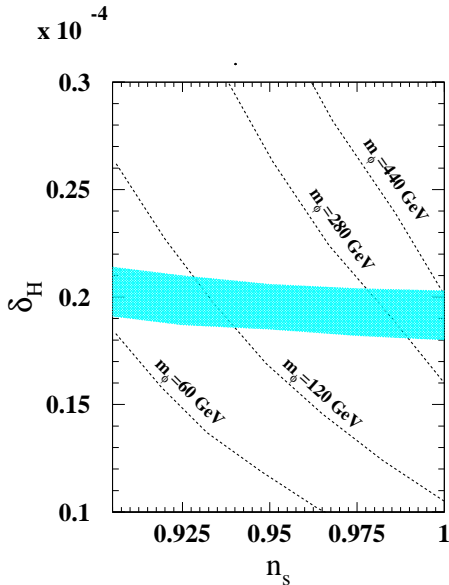


Figure 2: δ_H is plotted as a function of Δ^2 for different values of m_ϕ . We used $\lambda = 1$. The blue band denotes the experimentally allowed values of δ_H . Note that the parameter space allows the spectral tilt to be within $0.92 \leq n_s \leq 1.0$.

breaking sector and its communication with the observable sector may help to link the non-renormalizable A -term under consideration to the renormalizable ones. For instance in the Polonyi model the non-renormalizable A -term and the trilinear A -term can be related to each other: $A_6 = \frac{3-\sqrt{3}}{6-\sqrt{3}}A_3$.

SUGRA corrections often destroy the slow roll predictions of inflationary potentials; this is the notorious SUGRA- η problem [27]. In general, the effective potential depends on the Kähler potential K as $V \sim \left(e^{K(\varphi^*, \varphi)/M_P^2} V(\phi) \right)$ so that there is a generic SUGRA contribution to the flat direction potential of the type

$$V(\phi) = H^2 M_P^2 f\left(\frac{\phi}{M_P}\right), \quad (22)$$

where f is some function (typically a polynomial). Such a contribution usually gives rise to a Hubble induced correction to the mass of the flat direction with an unknown coefficient, which depends on the nature of the Kähler potential ⁶.

Let us compare the non-gravitational contribution, Eq. (2), to that of Hubble induced contribution, Eq. (22). Writing $f \sim (\phi/M_P)^p$ where $p \geq 1$ is some power,

⁶If the Kähler potential has a shift symmetry, then at tree level there is no Hubble induced correction. However, at one-loop level relatively small Hubble induced corrections can be induced [29, 30].

we see that non-gravitational part dominates whenever

$$H_{\text{inf}}^2 M_{\text{P}}^2 \left(\frac{\phi}{M_{\text{P}}} \right)^p \ll m_\phi^2 \phi_0^2, \quad (23)$$

so that the SUGRA corrections are negligible as long as $\phi_0 \ll M_{\text{P}}$, as is the case here (note that $H_{\text{inf}} M_{\text{P}} \sim m_\phi \phi_0$). The absence of SUGRA corrections is a generic property of this model. Note also that although non-trivial Kähler potentials give rise to non-canonical kinetic terms of squarks and sleptons, it is a trivial exercise to show that at sufficiently low scales, $H_{\text{inf}} \ll m_\phi$, and small VEVs, they can be rotated to a canonical form without affecting the potential ⁷.

After the end of inflation, the flat direction starts rolling towards its global minimum. At this stage the dominant term in the scalar potential will be: $m_\phi \phi^2/2$. Since the frequency of oscillations is $\omega \sim m_\phi \sim 10^3 H_{\text{inf}}$, the flat direction oscillates a large number of times within the first Hubble time after the end of inflation. Hence the effect of expansion is negligible.

We recall that the curvature of the potential along the angular direction is much larger than H_{inf}^2 . Therefore, the flat direction has settled at one of the minima along the angular direction during inflation from which it cannot be displaced by quantum fluctuations. This implies that no torque will be exerted, and hence the flat direction motion will be one dimensional, i.e. along the radial direction.

Flat direction oscillations excite those MSSM degrees of freedom which are coupled to it. The inflaton, either LLe or udd flat direction, is a linear combination of slepton or squark fields. Therefore inflaton has gauge couplings to the gauge/gaugino fields and Yukawa couplings to the Higgs/Higgsino fields. As we will see particles with a larger couplings are produced more copiously during inflaton oscillations. Therefore we focus on the production of gauge fields and gauginos. Keep in mind that the VEV of the MSSM flat direction breaks the gauge symmetry spontaneously, for instance **udd** breaks $SU(3)_C \times U(1)_Y$ while **LLe** breaks $SU(2)_W \times U(1)_Y$, therefore, induces a supersymmetry conserving mass $\sim g \langle \phi(t) \rangle$ to the gauge/gaugino fields in a similar way as the Higgs mechanism, where g is a gauge coupling. When the flat direction goes to its minimum, $\langle \phi(t) \rangle = 0$, the gauge symmetry is restored. In this respect the origin is a point of enhanced symmetry [31].

⁷The same reason, i.e. $H_{\text{inf}} \ll m_\phi$ also precludes any large Trans-Planckian correction. Any such correction would generically go as $(H_{\text{inf}}/M_*)^2 \ll 1$, where M_* is the scale at which one would expect Trans-Planckian effects to kick in.

There can be various phases of particle creation in this model, here we briefly summarize the most dominant one. Let us elucidate the physics, by considering the case when LLe flat direction is the inflaton ⁸.

An efficient bout of particle creation occurs when the inflaton crosses the origin, which happens twice in every oscillation. The reason is that fields which are coupled to the inflaton are massless near the point of enhanced symmetry. Mainly electroweak gauge fields and gauginos are then created as they have the largest coupling to the flat direction. The production takes place in a short interval, $\Delta t \sim (gm_\phi\phi_0)^{-1/2}$, where $\phi_0 \sim 10^{14}$ GeV is the initial amplitude of the inflaton oscillation, during which quanta with a physical momentum $k \lesssim (gm_\phi\phi_0)^{1/2}$ are produced. The number density of gauge/gaugino degrees of freedom is given by [35], see also [36]

$$n_g \approx \frac{(gm_\phi\phi_0)^{3/2}}{8\pi^3}. \quad (24)$$

As the inflaton VEV is rolling back to its maximum value ϕ_0 , the mass of the produced quanta $g\langle\phi(t)\rangle$ increases. The gauge and gaugino fields can (perturbatively) decay to the fields which are not coupled to the inflaton, for instance to (s)quarks. Note that (s)quarks are not coupled to the flat direction, hence they remain massless throughout the oscillations. The total decay rate of the gauge/gaugino fields is then given by $\Gamma = C(g^2/48\pi)g\phi$, where $C \sim \mathcal{O}(10)$ is a numerical factor counting for the multiplicity of final states.

The decay of the gauge/gauginos become efficient when [7]

$$\langle\phi\rangle \simeq \left(\frac{48\pi m_\phi\phi_0}{Cg^3}\right)^{1/2}. \quad (25)$$

Here we have used $\langle\phi(t)\rangle \approx \phi_0 m_\phi t$, which is valid when $m_\phi t \ll 1$, and $\Gamma \simeq t^{-1}$, where t represents the time that has elapsed from the moment that the inflaton crossed the origin. Note that the decay is very quick compared with the frequency of inflaton oscillations, i.e. $\Gamma \gg m_\phi$. It produces relativistic (s)quarks with an energy [7]:

$$E = \frac{1}{2}g\phi(t) \simeq \left(\frac{48\pi m_\phi\phi_0}{Cg}\right)^{1/2}. \quad (26)$$

⁸Reheating happens quickly due to a flat direction motion which is *strictly* one dimensional in our case. Our case is really exceptional, usually, the flat direction motion is restricted to a plane, which precludes preheating all together, for instance see [32, 33].

The ratio of energy density in relativistic particles thus produced ρ_{rel} with respect to the total energy density ρ_0 follows from Eqs. (24,26):

$$\frac{\rho_{rel}}{\rho_0} \sim 10^{-2} g, \quad (27)$$

where we have used $C \sim \mathcal{O}(10)$. This implies that a fraction $\sim \mathcal{O}(10^{-2})$ of the inflaton energy density is transferred into relativistic (s)quarks every time that the inflaton passes through the origin. This is so-called instant preheating mechanism [37]⁹. It is quite an efficient mechanism in our model as it can convert almost all of the energy density in the inflaton into radiation within a Hubble time (note that $H_{inf}^{-1} \sim 10^3 m_\phi^{-1}$).

A full thermal equilibrium is reached when *a) kinetic* and *b) chemical equilibrium* are established. The maximum (hypothetical) temperature attained by the plasma would be given by:

$$T_{max} \sim V^{1/4} \sim (m_\phi \phi_0)^{1/2} \geq 10^9 \text{ GeV}. \quad (28)$$

This temperature may be too high and could lead to thermal overproduction of gravitinos [41]. However the dominant source of gravitino production in a thermal bath is scattering which include an on-shell gluon or gluino leg. However there exists a natural solution to this problem and we showed that the final reheat temperature is actually well below Eq. (28), i.e. $T_R \ll T_{max}$.

One comment is in order before closing this subsection. The gravitinos can also be created non-perturbatively during inflaton oscillations, both of the helicity $\pm 3/2$ [42] and helicity $\pm 1/2$ states [43]. In models of high scale inflation (i.e. $H_{inf} \gg m_{3/2}$) helicity $\pm 1/2$ states can be produced very efficiently (and much more copiously than helicity $\pm 3/2$ states). At the time of production these states mainly consist of the inflatino (inflaton's superpartner). However these fermions also decay in the form of inflatino, which is coupled to matter with a strength which is equal to that of the inflaton. Therefore, they inevitably decay at a similar rate as that of inflaton, and hence pose no threat to primordial nucleosynthesis [44].

In the present case $m_\phi \sim m_{3/2} \gg H_{inf}$. Therefore low energy supersymmetry breaking is dominant during inflation, and hence helicity $\pm 1/2$ states of the gravitino are not related to the inflatino (which is a linear combination of leptons or

⁹In a favorable condition the flat direction VEV coupled very weakly to the flat direction inflaton could also enhance the perturbative decay rate of the inflaton [38].

quarks) at any moment of time. As a result helicity $\pm 1/2$ and $\pm 3/2$ states are excited equally, and their abundances are suppressed due to kinematical phase factor. Moreover there will be no dangerous gravitino production from perturbative decay of the inflaton quanta [39, 40, 45]. The reason is that the inflaton is not a *gauge singlet* and has gauge strength couplings to other MSSM fields. This makes the *inflaton* \rightarrow *inflatino* + *gravitino* decay mode totally irrelevant.

Let us briefly discuss the cold dark matter issue. Note that our model of inflation is embedded within MSSM, it is a bonus that the cold dark matter candidate comes out quite naturally once we assume the R-parity. In a simple toy model like mSUGRA there are only four parameters and one sign. These are m_0 (the universal scalar soft breaking mass at the GUT scale M_G); $m_{1/2}$ (the universal gaugino soft breaking mass at M_G); A_0 (the universal trilinear soft breaking mass at M_G)¹⁰; $\tan\beta = \langle H_2 \rangle / \langle H_1 \rangle$ at the electroweak scale (where H_2 gives rise to u quark masses and H_1 to d quark and lepton masses); and the sign of μ , the Higgs mixing parameter in the superpotential ($W_\mu = \mu H_1 H_2$). Unification of gauge couplings within supersymmetry suggests that $M_G \simeq 2 \times 10^{16}$ GeV. The model parameters are already significantly constrained by different experimental results. In subsequent plots we show that there exists an interesting overlap between the constraints from inflation and the CDM abundance [10].

We show that the mSUGRA parameter space in figures 3, for $\tan\beta = 10$ and 40 with the *udd* flat direction using $\lambda = 1$ ¹¹. In the figures, we show contours correspond to $n_s = 1$ for the maximum value of $\delta_H = 2.03 \times 10^{-5}$ (at 2σ level) and $n_s = 1.0, 0.98, 0.96$ for $\delta_H = 1.91 \times 10^{-5}$. The constraints on the parameter space arising from the inflation appearing to be consistent with the constraints arising from the dark matter content of the universe and other experimental results. We find that $\tan\beta$ needs to be smaller to allow for smaller values of $n_s < 1$. It is also interesting to note that the allowed region of m_ϕ , as required by the inflation data for $\lambda = 1$ lies in the stau-neutralino coannihilation region which requires smaller values of the SUSY particle masses. The SUSY particles in this parameter space are, therefore, within the reach of the LHC very quickly. The detection of the region at the LHC has been

¹⁰The relationship between the two A terms, the trilinear, A_0 and the non-renormalizable A term in Eq.(2) can be related to each other, however, that depends on the SUSY breaking sector. For a Polonyi model, they are given by: $A = (3 - \sqrt{3})/(6 - \sqrt{3})A_0$ [7].

¹¹We have a similar figure for the flat direction LLe which we do not show in this paper. All the figures are for *udd* flat direction as an inflaton.

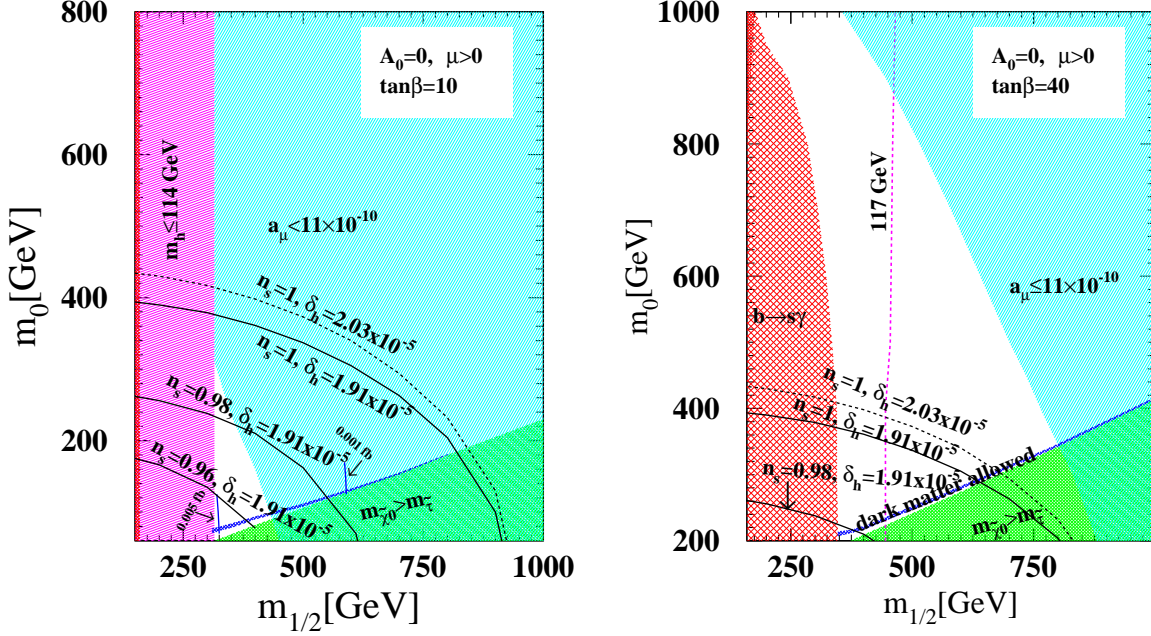


Figure 3: The contours for different values of n_s and δ_H are shown in the $m_0 - m_{1/2}$ plane for $\tan \beta = 10$ and $\tan \beta = 40$. We used $\lambda = 1$ for the contours. We show the dark matter allowed region narrow blue corridor, $(g-2)_\mu$ region (light blue) for $a_\mu \leq 11 \times 10^{-8}$, Higgs mass ≤ 114 GeV (pink region) and LEP II bounds on SUSY masses (red). We also show the dark matter detection rate by vertical blue lines.

considered in refs [46]. From the figures, one can also find that as $\tan \beta$ increases, the inflation data along with the dark matter, rare decay and Higgs mass constraint allow smaller ranges of $m_{1/2}$. For example, the allowed ranges of gluino masses are 765 GeV-2.1 TeV and 900 GeV-1.7 TeV for $\tan \beta = 10$ and 40 respectively [10].

So far we have chosen $\lambda = 1$. Now if λ is small e.g., $\lambda \lesssim 10^{-1}$, we find that the allowed values of m_ϕ to be large. In this case the dark matter allowed region requires the lightest neutralino to have larger Higgsino component in the mSUGRA model. As we will see shortly, this small value of λ is accommodated in $SO(10)$ type model. In figure 4, we show $n_s = 1, 0.98$ contours for $\delta_H = 1.91 \times 10^{-5}$ in the mSUGRA parameter space for $\tan \beta = 10$. In this figure, we find that n_s can not smaller than 0.97, but if we lower λ which will demand larger m_ϕ and therefore n_s can be lowered down to 0.92.

In the second panel of figure 4, we show the contours of λ for different values of m_ϕ which are allowed by n_s and $\delta_H = 1.91 \times 10^{-5}$. The blue bands show the

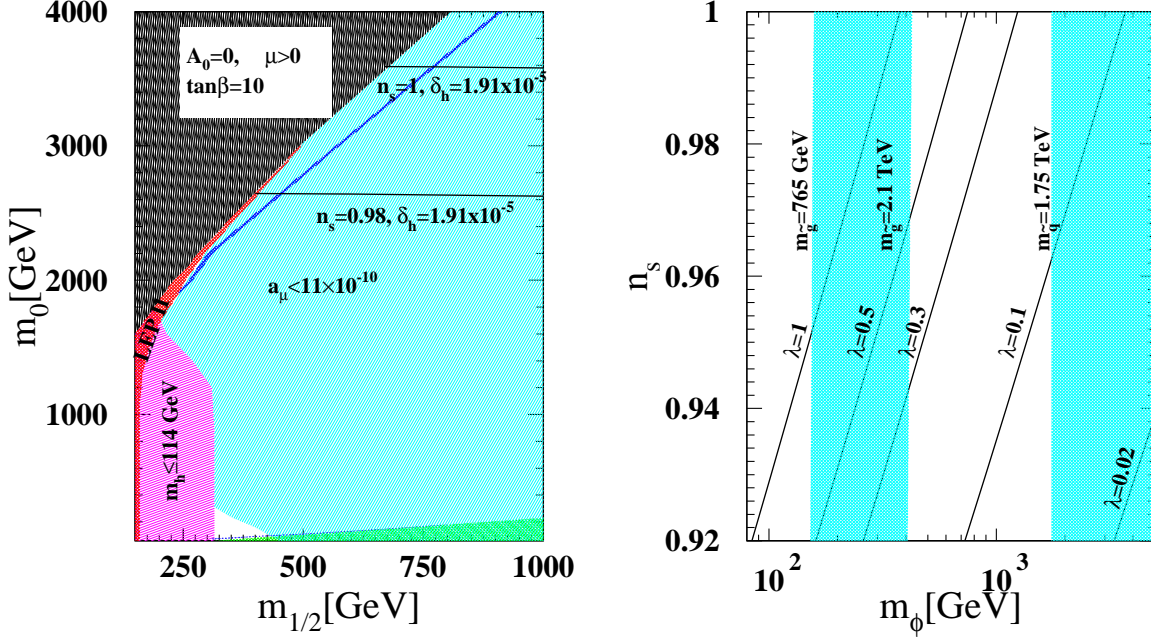


Figure 4: The contours for different values of n_s and δ_H are shown in the $m_0 - m_{1/2}$ plane for $\tan\beta = 10$. We used $\lambda = 0.1$ for the contours. We show the dark matter allowed region narrow blue corridor, g-2 region (light blue) for $a_\mu \leq 11 \times 10^{-8}$, Higgs mass ≤ 114 GeV (pink region) and LEP II bounds on SUSY masses (red). The black region is not allowed by radiative electroweak symmetry breaking. We use $m_t = 172.7$ GeV for this graph. In the right hand side plot we show the contours of λ for $\delta_H = 1.91 \times 10^{-5}$ in the n_s - m_ϕ plane. The blue band on the left is due to the stau-neutralino coannihilation region for $\tan\beta = 10$ and the blue band on the right (which continues beyond the plotting range) denotes the focus point region.

dark matter allowed regions for $\tan\beta = 10$. The band on the left is due to the stau-neutralino coannihilation region allowed by other constraints and the allowed values of λ are 0.3-1. The first two generation squarks masses are 690 GeV and 1.9 TeV for the minimum and maximum values of m_ϕ allowed by the dark matter and other constraints. The gluino masses for these are 765 GeV and 2.1 TeV respectively. The band is slightly curved due to the shifting of ϕ_0 as a function λ . (We solve for SUSY parameters from the inflaton mass at ϕ_0). The band on the right which continues beyond the plotting range of the figure ?? is due to the Higgsino dominated dark matter. We find that λ is mostly ≤ 0.1 in this region and $m_\phi > 1.9$ TeV. In this case the squark masses are much larger than the gluino mass since m_0 is much larger than $m_{1/2}$.

There have been other scintillating developments in embedding inflation, particularly, realizing saddle point inflation in particle physics. The flatness of the inflaton potential can be accounted for a weakness of the Dirac Yukawa couplings for the observed neutrino mass spectrum [8, 13]. In Ref. [13], we provided an example where part of the inflaton flat direction acts as a thermal dark matter candidate while part of it decays into the SM baryons. The detection of dark matter and the neutrino properties in neutrino-less double beta decay experiments would shed important lights on the inflaton origin.

To summarize, in near future it will be possible to unveil the origin of the inflaton in a terrestrial laboratory such as the LHC and the neutrino-less double beta decay experiments. Recognizing the inflaton dressed with a SM gauge group will be the most cherished success of an inflationary paradigm.

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